

Large $\tan\beta$ SUSY QCD corrections to $B \rightarrow X_s\gamma$ *

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Abstract

The charged-Higgs boson contributions to the Wilson coefficients C_7 and C_8 , relevant for the decay $B \rightarrow X_s\gamma$, are discussed in supersymmetric models at large $\tan\beta$. These contributions receive two-loop $\mathcal{O}(\alpha_s \tan\beta)$ corrections by squark-gluino subloops, which are possibly large and nondecoupling in the limit of heavy superpartners. In previous studies, the relevant two-loop Feynman integrals were approximated by using an effective two-Higgs-doublet lagrangian. However, this approximation is theoretically justified only when the typical supersymmetric scale M_{SUSY} is sufficiently larger than the electroweak scale $m_{\text{weak}} \sim (m_W, m_t)$ and the mass of the charged-Higgs boson m_{H^\pm} . Here we evaluate these two-loop integrals exactly and compare the results with the existing, approximated ones. We then examine the validity of this approximation beyond the region where it has been derived, i.e. for $m_H \gtrsim M_{\text{SUSY}}$ and/or $M_{\text{SUSY}} \sim m_{\text{weak}}$.

1 Introduction

The inclusive width of the radiative decays of the B mesons, $B \rightarrow X_s\gamma$, is well described by the short-distant processes $b \rightarrow s\gamma$ and $b \rightarrow sg$, since nonperturbative hadronic corrections are small and well under control. The partonic processes have been evaluated within the Standard Model (SM) up to the next-to-leading order in QCD [1] and partially beyond [2]. Because in the SM the processes $b \rightarrow s\gamma$ and $b \rightarrow sg$ occur through loops with W^\pm and top quark, possible new physics beyond the SM may contribute at the same level in perturbation. The rather good agreement between the SM prediction and recent experimental results [3] for the branching ratio $\text{BR}(B \rightarrow X_s\gamma)$, therefore, allows already to constrain some extensions of the SM.

In the minimal supersymmetric standard model (MSSM), new loop contributions to the decays $b \rightarrow s\gamma$ and $b \rightarrow sg$ come [4, 5] from the charged-Higgs boson H^\pm , charginos, gluino and neutralino. Their contributions are often comparable to or even larger than the SM one. For generic models, these new contributions have been calculated [6] at the leading-order precision in QCD. Higher-order QCD and SUSY QCD corrections to these contributions have been evaluated [7, 8, 9] for specific scenarios. One important finding is that the gluino may induce $\mathcal{O}(\alpha_s \tan\beta)$ corrections [8, 9] to these beyond-SM contributions. For models with very large $\tan\beta$, the ratio of two Higgs

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VEVs, these corrections can be comparable to the leading-order contributions and significantly affect the constraints [10] on the charged-Higgs boson and SUSY particles from the experiments.

Here we focus on the contribution of the charged-Higgs boson H^\pm in large- $\tan\beta$ scenarios and analyze the two-loop $\mathcal{O}(\alpha_s \tan\beta)$ corrections. In previous studies [8, 9], squarks and gluino are assumed to be sufficiently heavier than the electroweak scale and charged-Higgs boson. Under this restriction, the dominant part of these corrections has been evaluated by using an effective two-Higgs-doublet (2HD) lagrangian where squarks and gluino are integrated out. This approach gives rather compact and simple approximated formulas for the $\mathcal{O}(\alpha_s \tan\beta)$ corrections. However, the validity of this approximation is not theoretically justified beyond the parameter range treated in Refs. [8, 9], i.e. for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$ and/or $M_{\text{SUSY}} \sim m_{\text{weak}}$. It is important to examine, in such cases, how far the approximation in Refs. [8, 9] may deviate from the result of the exact two-loop Feynman integrals.

In this talk, we report on the calculation of the charged-Higgs boson contribution to the Wilson coefficients C_7 and C_8 , related to the processes $b \rightarrow s\gamma$ and $b \rightarrow sg$, to $\mathcal{O}(\alpha_s \tan\beta)$, by exact evaluation of the relevant two-loop diagrams. We first present the origin of the $\mathcal{O}(\alpha_s \tan\beta)$ corrections to the H^\pm contributions and list all necessary diagrams. We next review the approximation in Refs. [8, 9], here called the nondecoupling approximation. Finally we make a numerical comparison of the exact result and the nondecoupling approximation of the $\mathcal{O}(\alpha_s \tan\beta)$ results for the H^\pm contributions to $C_7(\mu_W)$ and $C_8(\mu_W)$, and discuss the validity of the approximation. A more complete discussion is presented in Refs. [11, 12].

2 $\mathcal{O}(\alpha_s \tan\beta)$ corrections to the H^\pm contribution

In the MSSM with large $\tan\beta$, the dominant part of the one-loop H^\pm contributions to the $b \rightarrow s\gamma$ and $b \rightarrow sg$ decays comes from the diagrams in Fig. 1, where the photon or gluon is to be attached to the t or the H^\pm lines. The enhancing factor $\tan\beta$ at the $\bar{t}_L b_R H^+$ vertex is cancelled by the suppressing factor $\cot\beta$ at the $\bar{s}_L t_R H^-$ vertex.

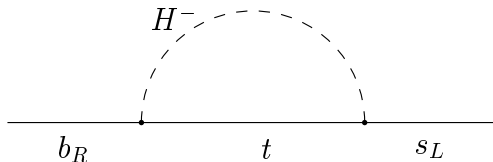


Figure 1: $b \rightarrow s\gamma$ and $b \rightarrow sg$ by the one-loop charged-Higgs boson exchange. The photon or gluon is to be attached at the t or H^- lines.

It has been shown [8, 9] that these H^\pm contributions receive $\mathcal{O}(\alpha_s \tan\beta)$ corrections from the squark-gluino subloops, which are potentially large for large $\tan\beta$. These corrections may arise from

1. counterterm for the $H^+ \bar{t}_L b_R$ coupling [13, 14, 15], coming from the mass corrections δm_b ;

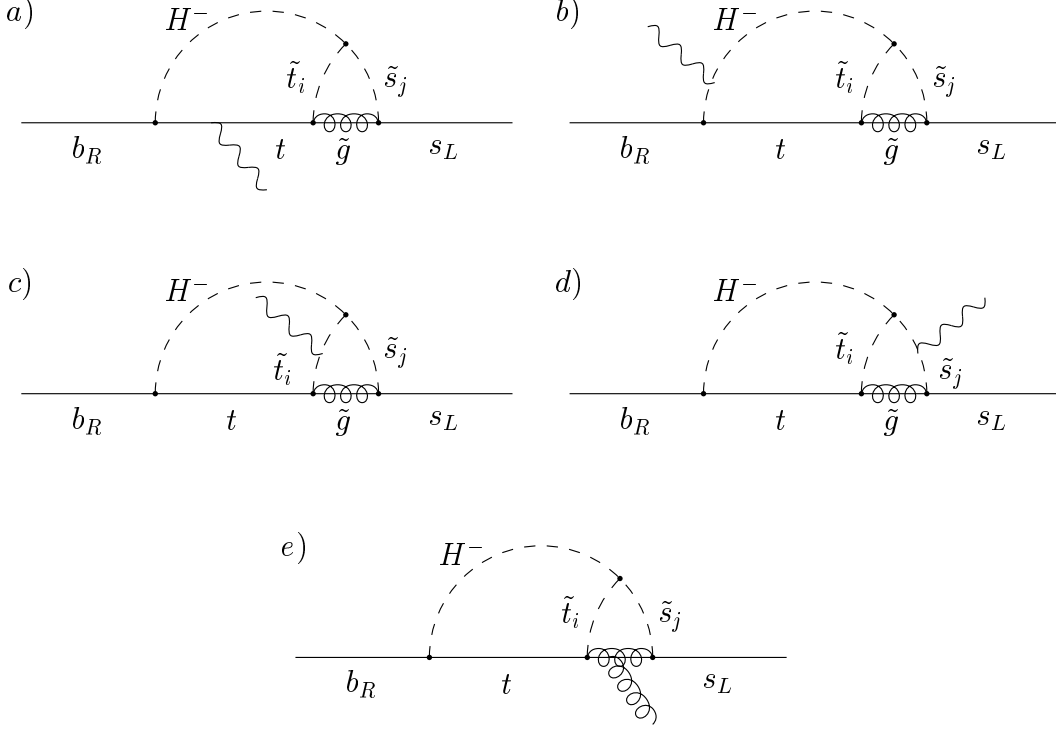


Figure 2: H^\pm mediated diagrams contributing at order $\mathcal{O}(\alpha_s \tan \beta)$ to the decays $b \rightarrow s\gamma$ and $b \rightarrow sg$. The photon must be replaced by a gluon and vice versa, whenever possible.

2. proper vertex corrections to the $H^- \bar{s}_L t_R$ coupling [8, 9];
3. effective four-point couplings $H^- \bar{s} t \gamma$ and $H^- \bar{s} t g$, as well as the $H^- \bar{s}_L t_L$ coupling, generated by squark-gluino subloops.

The two-loop diagrams relevant to the corrections of type 2 and 3, listed above, are shown in Fig. 2, where the photon must be replaced by a gluon, and vice versa, whenever possible. Note that, while the one-loop diagram in Fig. 1 has an chirality flip in the internal top quark line, the diagrams in Fig. 2 can have such a chirality flip also on the \tilde{t} -squark line, giving rise to the effective $H^- \bar{s}_L t_L$ coupling mentioned above. Both eigenstates of the \tilde{t} -squark, \tilde{t}_1 and \tilde{t}_2 , but only one \tilde{s} -squark eigenstate, the left-handed one, contribute in Fig. 2.

We calculate the H^\pm contributions to the Wilson coefficients C_7 and C_8 at the electroweak scale $\mu_W = m_W$, including the $\mathcal{O}(\alpha_s \tan \beta)$ corrections. Our normalization of $C_7(\mu_W)$ and $C_8(\mu_W)$ is the conventional one, as follows from the definition of the effective Hamiltonian,

$$H_{\text{eff}} \supset -\frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} (C_7(\mu) \mathcal{O}_7(\mu) + C_8(\mu) \mathcal{O}_8(\mu)) . \quad (1)$$

and of the operators \mathcal{O}_7 and \mathcal{O}_8 ,

$$\mathcal{O}_7(\mu) = \frac{e}{16\pi^2} m_b(\mu) \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu} , \quad \mathcal{O}_8(\mu) = \frac{g_s}{16\pi^2} m_b(\mu) \bar{s}_L \sigma^{\mu\nu} T^a b_R G_{\mu\nu}^a , \quad (2)$$

where $F_{\mu\nu}$ and $G_{\mu\nu}^a$ are the field strengths of the photon and the gluon, respectively.

We denote by $C_{7,H}(\mu_W)$ and $C_{8,H}(\mu_W)$ the $\tan\beta$ -unsuppressed H^\pm contribution to $C_7(\mu_W)$ and $C_8(\mu_W)$, respectively. They are decomposed as

$$C_{i,H}(\mu_W) = \frac{1}{1 + \Delta_{b_R,b} \tan\beta} \left[C_{i,H}^0(\mu_W) + \Delta C_{i,H}^1(\mu_W) \right], \quad (3)$$

where $C_{i,H}^0(\mu_W)$ and $\Delta C_{i,H}^1(\mu_W)$ are the contributions of the one-loop diagram in Fig. 1 and the two-loop diagrams in Fig. 2, respectively. The overall factor $1/(1 + \Delta_{b_R,b} \tan\beta)$ represents the correction to the $H^+ \bar{t}_L b_R$ Yukawa coupling coming from the correction to m_b [13, 14]. The one-loop function $\Delta_{b_R,b}$ is given in Ref. [12] and of the order of $\alpha_s \mu m_{\tilde{g}} / M_{\text{SUSY}}^2 \sim \alpha_s M_{\text{SUSY}}^0$. In the large M_{SUSY} limit, the contributions from $\Delta_{b_R,b}$ and the vertex corrections to the $H^- \bar{s}_L t_R$ coupling are nondecoupling, while all other contributions of Fig. 2 are decoupling.

We calculate all contributions of the two-loop diagrams in Fig. 2 by exact evaluation of the loop integrals, making use of results and techniques in Ref. [16]. The explicit forms of the Feynman integrals for these diagrams are listed in Ref. [12].

3 Nondecoupling approximation vs. exact calculation

In Refs. [8, 9], the calculations of the $\mathcal{O}(\alpha_s \tan\beta)$ corrections were performed under the assumption that all squarks and gluino, around M_{SUSY} , are sufficiently heavier than the top quark and the W boson, whereas m_{H^\pm} is around the electroweak scale $m_{\text{weak}} \sim m_W, m_t$. The squark-gluino subloop corrections to the H^\pm contributions were described in terms of an effective 2HD lagrangian, in which squarks and gluino are integrated out. In the following we call the approximation in Refs. [8, 9] the nondecoupling approximation, since it collects all nondecoupling parts of the $\mathcal{O}(\alpha_s \tan\beta)$ corrections. Strictly speaking, however, it includes parts of the formally decoupling $\mathcal{O}(m_{\text{weak}}^2 / M_{\text{SUSY}}^2)$ contributions through the masses and couplings of squarks [17]. For the contribution from δm_b , the use of the effective 2HD lagrangian allows us to resum higher-order $\mathcal{O}((\alpha_s \tan\beta)^n)$ terms [15] in Eq. (3), by putting $\Delta_{b_R,b}$ in the denominator.

For $\Delta C_{i,H}^1(\mu_W)$ coming from the diagrams in Fig. 2, the nondecoupling approximation is obtained by retaining only the diagrams a) and b), with chirality flip on the t -quark line only, and evaluating the squark-gluino subloops at vanishing external momenta. By this approximation, the original two-loop Feynman integrals for the $\mathcal{O}(\alpha_s \tan\beta)$ corrections are factorized into two one-loop diagrams, taking rather compact forms.

We show one example to illustrate the difference between the nondecoupling approximation and the exact calculation. The contribution of the diagram Fig. 2a), with chirality flip on the t -quark line, is proportional to the integral

$$\begin{aligned} I_{ti2} &= \mu m_{\tilde{g}} \int \frac{d^4 k}{(2\pi)^4} \int \frac{d^4 l}{(2\pi)^4} \frac{1}{[k^2 - m_{\tilde{t}}^2]^2 [k^2 - m_{H^\pm}^2] [(l+k)^2 - m_{\tilde{t}}^2] [l^2 - m_{\tilde{s}}^2] [l^2 - m_{\tilde{g}}^2]} \\ &\quad \times \left\{ \frac{l \cdot k}{l^2 - m_{\tilde{g}}^2} - \frac{2k^2}{k^2 - m_{\tilde{t}}^2} \right\}. \end{aligned} \quad (4)$$

The loop momenta k and l represent the momenta of (t, H^\pm) and SUSY particles, respectively. In the nondecoupling approximation, the k -dependence of the \tilde{t}_i line is neglected, i.e. $(l+k)^2 - m_{\tilde{t}_i}^2$ is

replaced by $l^2 - m_{\tilde{t}_i}^2$. The term proportional to $l \cdot k$ is then dropped and the integral is factorized into two one-loop integrals as

$$I_{ti2}|_{\text{nondec}} = \mu m_{\tilde{g}} \int \frac{d^4 k}{(2\pi)^4} \frac{-2k^2}{[k^2 - m_t^2]^3 [k^2 - m_{H^\pm}^2]} \int \frac{d^4 l}{(2\pi)^4} \frac{1}{[l^2 - m_{\tilde{t}_i}^2] [l^2 - m_s^2] [l^2 - m_g^2]}. \quad (5)$$

In the nondecoupling approximation, $\mathcal{O}((m_{\text{weak}}^2, m_{H^\pm}^2)/M_{\text{SUSY}}^2)$ terms which may come from the k -dependence of the squark-gluino subloops of the diagrams in Fig. 2a,b), as well as the whole contributions of the diagrams in Fig. 2c-e), are neglected. The resulting deviation of this approximation from the exact two-loop calculation is, therefore, expected to become large when M_{SUSY} is not much heavier than m_{weak} and/or m_{H^\pm} .

Since the condition for the theoretical justification of the nondecoupling approximation, $m_{\text{weak}}^2 \sim m_{H^\pm}^2 \ll M_{\text{SUSY}}^2$, is often violated in well-known scenarios for the SUSY breaking mechanism, it is very important to study how far this approximation may be applied beyond this restricted parameter region. Clearly, a definite answer to this question will be given by the exact calculation of all the two-loop diagrams in Fig. 2, without any assumption on the relative size of m_{H^\pm} , M_{SUSY} , and m_{weak} .

4 Numerical results

We present numerical results for the H^\pm contributions, $C_{i,H}(\mu_W)$ ($i = 7, 8$) shown in Eq. (3), at the scale $\mu_W = M_W$. We make a comparison of the results $C_{i,H}(\mu_W)|_{\text{exact}}$ obtained from the exact two-loop integrals with the nondecoupling approximation $C_{i,H}(\mu_W)|_{\text{nondec}}$.

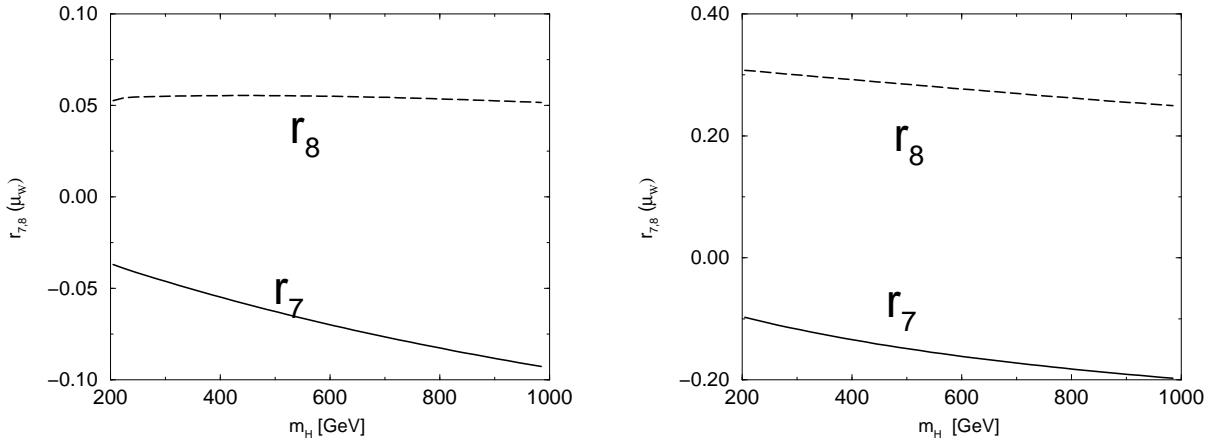


Figure 3: Relative deviations $r_i(\mu_W)$ ($i = 7, 8$) of the nondecoupling approximations of the $\mathcal{O}(\alpha_s \tan \beta)$ Wilson coefficients $C_{i,H}(\mu_W)$ from the exact two-loop results, as defined in the text, for the spectrum I (left) and spectrum II (right).

In Fig. 3, we plot the ratios

$$r_i(\mu_W) \equiv \frac{C_{i,H}(\mu_W)|_{\text{nondec}} - C_{i,H}(\mu_W)|_{\text{exact}}}{C_{i,H}(\mu_W)|_{\text{exact}}} \quad (i = 7, 8), \quad (6)$$

showing the relative deviation of the nondecoupling approximation from the exact two-loop calculation, as functions of m_{H^\pm} . The correction $\Delta_{bR,b}$ in Eq. (3), coming from the mass correction δm_b , cancel out in the ratios r_i .

Two sets of the parameters for squarks and gluino are used in Fig. 3. For an example of a heavier SUSY spectrum, called here spectrum I, we have chosen $(m_{\tilde{s}_L}, m_{\tilde{t}_1}, m_{\tilde{t}_2}) = (700, 500, 450)$ GeV, the left-right mixing angle of \tilde{t} -squarks $\cos \theta_t = 0.8$, $m_{\tilde{g}} = 600$ GeV, and $\mu = 550$ GeV. For a lighter SUSY spectrum, spectrum II, we set $(m_{\tilde{s}_L}, m_{\tilde{t}_1}, m_{\tilde{t}_2}) = (350, 400, 320)$ GeV, $\cos \theta_t = 0.8$, $m_{\tilde{g}} = 300$ GeV, and $\mu = 450$ GeV. As for other input parameters, we have used $\tan \beta = 30$, $m_t(\mu_W) = 176.5$ GeV, which corresponds to a pole mass $M_t = 175$ GeV, and $\alpha_s(\mu_W) = 0.12$.

For the spectrum I, the difference between the exact calculation and the nondecoupling approximation is very small in the whole range of m_{H^\pm} , even for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$. This is a surprising result since, as discussed in Sect. 3, the $\mathcal{O}(m_{H^\pm}^2/M_{\text{SUSY}}^2)$ deviation was expected to be large in this region. In the case of the spectrum II, $r_{7,8}$ become larger. The corrections beyond the nondecoupling approximation are of the same order of the $\text{SU}(2) \times \text{U}(1)$ breaking effects in the SUSY particle subloops [17] and are no longer negligible. Nevertheless, r_7 and r_8 remain of the same order of magnitude for increasing m_{H^\pm} , up to $m_{H^\pm} \gg M_{\text{SUSY}}$. In both cases, the main part of the difference between the result of the nondecoupling approximation and the exact two-loop result comes from the diagram in Fig. 2a) and, for $C_{8,H}$, also from the diagram in Fig. 2e).

To understand the results for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$ qualitatively, we again consider the diagram in Fig. 2a), with chirality flip on the top quark line. When m_{H^\pm} is sufficiently larger than m_t , this diagram gives the largest contribution to $\Delta C_{7,H}^1(\mu_W)$ and $\Delta C_{8,H}^1(\mu_W)$. It is proportional to the integral I_{ti2} in Eq. (4). For the following discussion we rewrite I_{ti2} in the form

$$I_{ti2}(m_t, m_{H^\pm}, m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}) = \int \frac{d^4 k}{(2\pi)^4} \frac{k^2}{[k^2 - m_t^2]^3 [k^2 - m_{H^\pm}^2]} Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}). \quad (7)$$

$Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}})$ represents the form factor for the effective vertex $H^- \bar{s}_L t_R$ generated by the squark-gluino loops and is given by:

$$Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}) = \mu m_{\tilde{g}} \left[-2F + (k^2 - m_t^2)G \right] (k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2), \quad (8)$$

with

$$F(k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2) = \int \frac{d^4 l}{(2\pi)^4} \frac{1}{[(l+k)^2 - m_{\tilde{t}_i}^2] [l^2 - m_{\tilde{s}}^2] [l^2 - m_{\tilde{g}}^2]}, \quad (9)$$

$$k^\mu G(k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2) = \int \frac{d^4 l}{(2\pi)^4} \frac{l^\mu}{[(l+k)^2 - m_{\tilde{t}_i}^2] [l^2 - m_{\tilde{s}}^2] [l^2 - m_{\tilde{g}}^2]^2}. \quad (10)$$

The nondecoupling approximation of I_{ti2} , shown in Eq. (5), is obtained by replacing Y_{ti2} in Eq. (7) with

$$Y_{ti2}|_{\text{nondec}} = -2\mu m_{\tilde{g}} F(0; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2), \quad (11)$$

which is an $\mathcal{O}(M_{\text{SUSY}}^0)$ constant with respect to k^2 . To simplify our discussion, we set hereafter $(m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}, \mu)$ equal to M_{SUSY} .

For $|k^2| \ll M_{\text{SUSY}}^2$, $F(k^2; M_{\text{SUSY}}^2)$ and $G(k^2; M_{\text{SUSY}}^2)$ behave as

$$\begin{aligned} F(k^2; M_{\text{SUSY}}^2) &= \mathcal{O}\left(\frac{1}{M_{\text{SUSY}}^2}\right) + \mathcal{O}\left(\frac{k^2}{M_{\text{SUSY}}^4}\right), \\ G(k^2; M_{\text{SUSY}}^2) &= \mathcal{O}\left(\frac{1}{M_{\text{SUSY}}^4}\right). \end{aligned} \quad (12)$$

For $|k^2| \gg M_{\text{SUSY}}^2$, it is:

$$\begin{aligned} F(k^2; M_{\text{SUSY}}^2) &\rightarrow \mathcal{O}\left(\frac{1}{k^2} \ln \frac{k^2}{M_{\text{SUSY}}^2}\right), \\ G(k^2; M_{\text{SUSY}}^2) &\rightarrow \mathcal{O}\left(\frac{1}{k^4} \ln \frac{k^2}{M_{\text{SUSY}}^2}\right). \end{aligned} \quad (13)$$

The behavior of $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ is therefore

$$Y_{ti2}(k^2; M_{\text{SUSY}}^2) \rightarrow \begin{cases} Y_{ti2}|_{\text{nondec}} + \mathcal{O}\left(\frac{k^2}{M_{\text{SUSY}}^2}, \frac{m_t^2}{M_{\text{SUSY}}^2}\right) & (|k^2| \ll M_{\text{SUSY}}^2), \\ \mathcal{O}\left(\frac{M_{\text{SUSY}}^2}{k^2} \ln \frac{k^2}{M_{\text{SUSY}}^2}\right) & (|k^2| \gg M_{\text{SUSY}}^2), \end{cases} \quad (14)$$

which supports the naive expectation that a substantial deviation of $I_{ti2}(m_t, m_{H^\pm}, M_{\text{SUSY}}^2)$ from its nondecoupling approximation $I_{ti2}(m_t, m_{H^\pm}, M_{\text{SUSY}}^2)|_{\text{nondec}}$ may arise for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$.

However, the factor multiplying $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ in Eqs. (7) plays an important role, leading to the fact that this expectation does not hold in the case in which M_{SUSY} is not rather light. Since for $|k^2| \gg m_{H^\pm}^2$ this factor drops as $d^4 k/k^6$, the integral (7) gets its largest contribution from the region $|k^2| \lesssim m_{H^\pm}^2$. A closer inspection actually shows that it is the region of small $|k^2|$, up to $|k^2| = \mathcal{O}(m_t^2)$, which determines the bulk of the value of this integral. If M_{SUSY} is sufficiently larger than m_t , $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ does not deviate substantially from $Y_{ti2}|_{\text{nondec}}$ in this region. This explains the smallness of the deviation for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$ shown in Fig. 3.

5 Conclusion

We have studied the $\mathcal{O}(\alpha_s \tan \beta)$ corrections to the H^\pm contributions to the Wilson coefficients relevant for the decay $B \rightarrow X_s \gamma$, in the MSSM with large $\tan \beta$. These corrections are generated by the shift of the b -quark mass in the Higgs-quark couplings and by the dressing of the one-loop H^\pm diagrams with squark-gluino subloops, as shown in Fig. 2.

In this talk, we have focused on the latter class of corrections. In previous studies [8, 9], the contributions from these two-loop diagrams were calculated in an approximated way, by using an effective 2HD lagrangian formalism in which squarks and gluino are integrated out. This method, here called the nondecoupling approximation, is theoretically justified in the case $m_{\text{weak}}^2 \sim m_{H^\pm}^2 \ll M_{\text{SUSY}}^2$, and gives rather compact forms for these corrections. However, the deviation from the exact two-loop result was, in principle, expected to be of $\mathcal{O}(m_{\text{weak}}^2, m_{H^\pm}^2/M_{\text{SUSY}}^2)$, and to become significant when $m_{\text{weak}} \leq M_{\text{SUSY}}$ and/or $m_{H^\pm} \geq M_{\text{SUSY}}$.

We have calculated the contributions of the two-loop diagrams in Fig. 2 exactly, without assuming any patterns for the mass of the particle involved, and compared the results with the non-decoupling approximation. Surprisingly, the difference between the nondecoupling approximation in Refs. [8, 9] and the exact two-loop result was shown to be quite small, even for $m_{H^\pm} \gtrsim M_{\text{SUSY}}$, provided M_{SUSY} is sufficiently larger than m_{weak} . The unexpected absence of large deviation for the case of $m_{H^\pm} \gtrsim M_{\text{SUSY}}$ with $M_{\text{SUSY}}^2 \gg m_{\text{weak}}^2$ can be understood from the structure of the relevant two-loop integrals. In contrast, nonnegligible deviation appeared for M_{SUSY} not much larger than m_{weak} .

We have illustrated our findings by showing the H^\pm contributions to the Wilson coefficients $C_{7,H}$ and $C_{8,H}$ at the electroweak matching scale μ_W , for different spectra of the gluino, squarks and H^\pm . Analyses of C_7 and C_8 at a low scale $\sim m_b$, including other contributions than the H^\pm -mediated one, and of the actual branching ratio $\text{BR}(B \rightarrow X_s \gamma)$, will be presented in future work.

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